Hints of Standard Model Higgs Boson at the LHC and Light Dark Matter Searches

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Abstract

The most recent results of searches at the LHC for the Higgs boson h have turned up possible hints of such a particle with mass m_h about 125 GeV consistent with standard model (SM) expectations. This has many potential implications for the SM and beyond. We consider some of them in the contexts of a simple Higgs-portal dark matter (DM) model, the SM plus a real gauge-singlet scalar field D as the DM candidate, and a couple of its variations. In the simplest model with one Higgs doublet and three or four generations of fermions, for D mass $m_D < m_h/2$ the invisible decay $h \to DD$ tends to have a substantial branching ratio. If future LHC data confirm the preliminary Higgs indications, m_D will have to exceed $m_h/2$. To keep the DM lighter than $m_h/2$, one will need to extend the model and also satisfy constraints from DM direct searches. The latter can be accommodated if the model provides sizable isospin violation in the DM-nucleon interactions. We explore this in a two-Higgs-doublet model combined with the scalar field D. This model can offer a 125-GeV SM-like Higgs and a light DM candidate having isospin-violating interactions with nucleons at roughly the required level, albeit with some degree of fine-tuning.

I. INTRODUCTION

The latest searches for the standard model (SM) Higgs boson performed by the ATLAS and CMS Collaborations at the LHC have come up with tantalizing hints of the particle [1]. They observed modest excesses of events compatible with a SM Higgs h having mass m_h in the range from 124 to 126 GeV, but the statistical significance of the excesses is not enough for making any conclusive statement on the Higgs existence or nonexistence [1]. Interestingly, these numbers are consistent with $m_h = 125^{+8}_{-10}$ GeV from the SM complete fit to electroweak precision data plus constraints from direct Higgs searches at LEP and Tevatron [2]. Needless to say, if upcoming measurements confirm these preliminary findings at the LHC to be glimpses of the SM Higgs, or a SM-like Higgs, the implications will be far-reaching for efforts to extend the SM, as physics beyond the SM is still necessary to account for, among other things, the observed evidence for dark matter. Particularly, all new-physics models would have to include such a spinless boson as one of their ingredients.

Simultaneously with the quest for the Higgs, a number of underground experiments have for years been searching directly for weakly interacting massive particle (WIMP) dark matter (DM) by looking for the recoil energy of nuclei caused by WIMPs colliding with nucleons. Intriguingly, some of these searches have turned up excess events which may have been WIMP signals. Specifically, the DAMA, CoGeNT, and CRESST-II Collaborations [3–5] have acquired data that seem to be pointing to light WIMPs of mass in the region roughly from 5 to 30 GeV and spin-independent WIMP-nucleon scattering cross-sections of order 10^{-42} to 10^{-40} cm^2, although the respective ranges preferred by these experiments do not fully agree with each other. In contrast, other direct searches, especially by the CDMS, XENON, and SIMPLE Collaborations [6–9], still have not produced any WIMP evidence. Although presently for WIMP masses under 15 GeV the latter null results are still controversial [10] and future WIMP searches with improved sensitivity may eventually settle this issue definitively, there may be alternative explanations worth exploring which can reconcile these disagreeing findings on DM.

Since the various DM searches employed different target materials for WIMP detection, one of the possible resolutions to the light-DM controversy that have been proposed is to allow large isospin violation in the WIMP-nucleon interactions [11–13]. It turns out that the tension between the conflicting direct-search results can be partially eased if the effective WIMP couplings f_p and f_n to the proton and neutron, respectively, satisfy the ratio $f_n/f_p \simeq -0.7$ [12–14].

In view of these developments in the hunts for the Higgs and for DM, it is of interest to consider some of their implications within the context of a relatively simple framework. For if the two sectors are intimately connected, detecting the signs of one of them could shine light on still hidden elements of the other.

The most economical model possessing both a Higgs boson and a WIMP candidate is the SM+D, which is the SM expanded with the addition of a real gauge-singlet scalar field D dubbed darkon acting as the WIMP [15–19]. This model predicts that the Higgs decay for $m_h \sim 120$-130 GeV is highly dominated by the invisible mode $h \to DD$ if the darkon mass $m_D < m_h/2$, except when m_D is close to, but less than, $m_h/2$ [17–19]. Thus, if a Higgs with $m_h \sim 125$ GeV and characteristics within the SM preference manifests itself unambiguously in LHC data, the parameter space of the SM+D with a light darkon will be strongly diminished. On the other hand, for $m_D \geq m_h/2$ the model is consistent with the existence of such a Higgs, as its decay pattern is unaffected at leading order by the presence of the darkon, although limited portions of this m_D zone are already excluded by direct-search data.

In order to have a Higgs compatible with the LHC indications as well as a light-WIMP candidate, one must therefore expand the SM+D. This motivates us to study in this paper a slightly extended model we call THDM+D, which is a two-Higgs-doublet model combined with a darkon.[1] It can offer such a Higgs and an ample amount of viable parameter space for the darkon. The model can also supply isospin-violating WIMP-nucleon interactions at about the desired level, although this will require fine-tuning to some extent. In addition, the potential presence of tree-level flavor-changing couplings of the neutral Higgs bosons in the THDM+D implies that the Higgs-mediated top-quark decays $t \to (u,c)DD$, if kinematically allowed, would contribute to the decays $t \to (u,c)$ plus missing energy. If these decays have rates that are sufficiently amplified to be measurable at the LHC, they can be another avenue to probe the darkon [20].

In the next section, we consider in more detail some implications of the possible discovery of a Higgs having SM-like properties for the SM+D with either three or four sequential generations of fermions (hereafter referred to as SM3+D or SM4+D, respectively). In Sec. III we describe the THDM+D with a SM-like Higgs and study its prediction for WIMP-nucleon cross-sections in the limit that isospin violation is negligible in the WIMP-nucleon effective couplings. In Sec. IV we treat the THDM+D case with isospin-violating WIMP-nucleon interactions. We give our conclusions in Sec. V.

II. STANDARD MODEL PLUS DARKON

In the DM sector of the SM+D, to ensure the stability of the darkon D as the DM, one assumes it to be a singlet under the gauge groups of the model and introduces a Z_2 symmetry under which $D \to -D$, all the other fields being unaffected. Requiring that the darkon Lagrangian be renormalizable then implies that it has the form [15, 16]

$$\mathcal{L}_D = \tfrac{1}{2}\partial^\mu D\,\partial_\mu D - \tfrac{1}{4}\lambda_D D^4 - \tfrac{1}{2}m_0^2 D^2 - \lambda D^2 H^\dagger H \ , \tag{1}$$

where λ_D, m_0, and λ are free parameters and H is the Higgs doublet containing the physical Higgs field h, in the notation of Ref. [17] which has additional details on the model. Clearly its darkon sector has very few free parameters, only two of which, besides m_h, are relevant here: the Higgs-darkon coupling λ, which determines the darkon relic density, and the darkon mass $m_D = (m_0^2 + \lambda v^2)^{1/2}$, where $v = 246\,\mathrm{GeV}$ is the Higgs vacuum expectation value (VEV).

As explained recently in Ref. [18], constraints on the darkon in the SM+D (either SM3+D or SM4+D) from a number of rare meson decays with missing energy and from DM direct searches, pending a definitive resolution to the light-WIMP puzzle, allow the darkon mass values $2.5\,\mathrm{GeV} \le m_D \le 15\,\mathrm{GeV}$ consistent with the light-WIMP hypothesis and those not far from, but smaller than, $m_h/2$. For such masses, the invisible decay mode $h \to DD$ can dominate the Higgs total width depending on m_h, making the Higgs still hidden from detection, and as a consequence significant portions of the m_h ranges in the SM3 (SM4) excluded by the current LHC data may be made viable again in the SM3+D (SM4+D).[2]

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[1] Various aspects of the THDM+D were previously addressed in Refs. [19–22].

[2] Here, as in Refs. [17, 18], all the fourth-generation fermions in the SM4+D are assumed to be unstable. If the fourth-generation neutrino is stable, it can be a DM candidate which may also contribute to the Higgs invisible decay and makes up a fraction of the DM relic density, as was studied in Ref. [23] in the absence of the darkon.

With the appearance of possible clues of a SM-like Higgs in the latest LHC data [1], we consider in this section some of the implications for the SM3+D and SM4+D.[3] We follow Ref. [18] to apply the procedure given in Ref. [17] for $3\,\text{GeV} \le m_D \le 400\,\text{GeV}$, but now with the specific selection $m_h = 125\,\text{GeV}$ for definiteness, in order to extract the darkon-Higgs coupling λ from the measured DM relic density $\Omega_D h^2 = 0.1123 \pm 0.0035$ [26]. We present the results for the SM3+D in Fig. 1(a), where the band width reflects the 90% confidence-level (CL) range $0.092 \le \Omega_D h^2 \le 0.118$ and the black-dotted sections are ruled out by direct-search limits. The λ values in the SM4+D (not drawn) are roughly similar and mostly somewhat lower than their SM3+D counterparts, by no more than ∼20% [17, 18]. In Fig. 1(b) we show

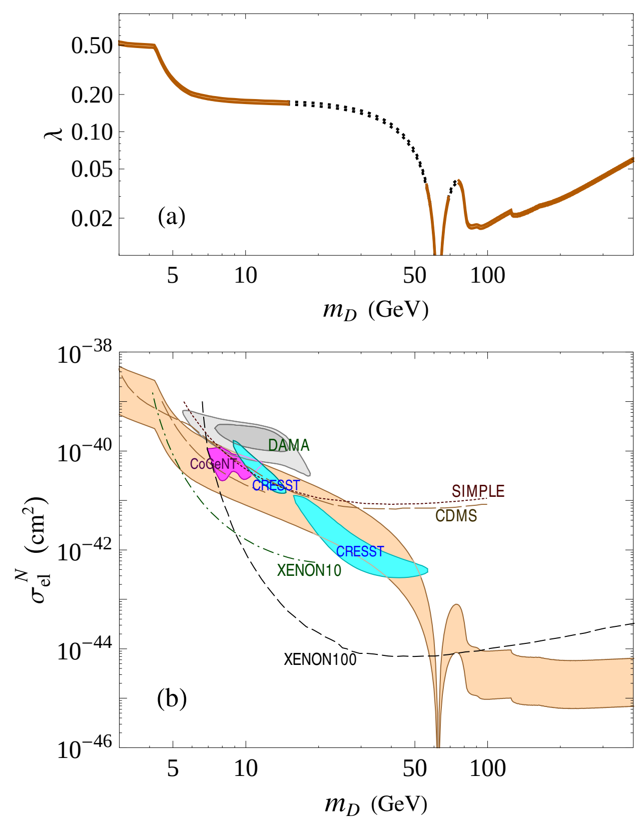

FIG. 1: (a) Darkon-Higgs coupling λ as a function of darkon mass m_D for Higgs mass $m_h = 125\,\text{GeV}$ in SM3+D. (b) The corresponding darkon-nucleon cross-section $\sigma_{\rm el}^N$, compared to 90%-CL upper limits from CDMS (brown long-dashed curves) [6], XENON10 (green dot-dashed curve) [7], XENON100 (black short-dashed curve) [8], and Stage 2 of SIMPLE (dark-red dotted curve) [9], as well as the 90%-CL (magenta) signal region suggested by CoGeNT [4], a gray (lighter gray) region compatible with the DAMA modulation signal at the 3σ (5σ) level [3, 27], and two 2σ-confidence (cyan) areas representing the CRESST-II result [5]. The black-dotted parts of the curve in (a) are disallowed by the direct-search limits in (b), after Ref. [18].

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[3] Some other aspects of the darkon model or its close variants and its potential impact on Higgs searches were treated before in Ref. [24]. Alternative scenarios involving scalar dark matter which may also affect the Higgs sector were dealt with recently in Ref. [25].

the darkon-nucleon elastic cross-section $\sigma_{\rm el}^N$ computed using the parameter choices in Fig. 1(a). The band width of the $\sigma_{\rm el}^N$ curve arises mainly from the sizable uncertainty of the Higgs-nucleon coupling, $0.0011 \le g_{NNh} \le 0.0032$ [18], due to its dependence on the pion-nucleon sigma term $\sigma_{\pi N}$ which is not well-determined [28]. Although the displayed $\sigma_{\rm el}^N$ at each m_D value now varies by up to an order of magnitude, we get a more realistic picture of how the model confronts the latest data from the leading direct-searches for WIMP DM, which are also shown in Fig. 1(b). In the SM4+D the majority of the $\sigma_{\rm el}^N$ numbers are $\sim$50% higher than their SM3+D counterparts [17].

If $m_h > 2m_D$, the invisible decay channel $h \to DD$ will become open with branching ratio $\mathcal{B}(h \to DD) = \Gamma(h \to DD)/\Gamma_h^{\rm SM+D}$, where $\Gamma(h \to DD) = \lambda^2 v^2 \big(1 - 4m_D^2/m_h^2\big)^{1/2}/(8\pi m_h)$ and $\Gamma_h^{\rm SM+D} = \Gamma_h^{\rm SM} + \Gamma(h \to DD)$ includes the Higgs total width $\Gamma_h^{\rm SM}$ in the SM3 or SM4 without the darkon. From the λ values obtained above, we plot $\mathcal{B}(h \to DD)$ in Fig. 2(a), where the dotted portions are again excluded. In Fig. 2(b) we display the corresponding reduction factor [16, 18]

$$\mathcal{R} \;=\; \frac{\mathcal{B}\big(h \to X\bar{X}\big)}{\mathcal{B}\big(h \to X\bar{X}\big)_{\rm SM}} \;=\; \frac{\Gamma_h^{\rm SM}}{\Gamma_h^{\rm SM} + \Gamma(h \to DD)_{\rm SM+D}} \tag{2}$$

which lowers all the SM3 (SM4) Higgs branching ratios in the SM3+D (SM4+D) in the same way. It is clear from these graphs that for $m_D \le 15\,{\rm GeV}$ the SM3+D Higgs would be mainly invisible. Consequently the observation of a 125-GeV SM3-like Higgs having nonnegligible branching ratios in the visible channels would imply the exclusion of the SM3+D light-darkon region. Moreover, the only surviving masses would be restricted to the vicinity of, but below, the boundary $m_D = m_h/2$. A similar conclusion was drawn in Ref. [29].

In the SM4+D, the effect of a light darkon, via the reduction factor $\mathcal{R}$, on the Higgs production rate may be ameliorated by the enhancement of the gluon-fusion cross-section $\sigma(gg \to h)$ by up to $\sim$9 times due to the fourth-generation quarks [30]. To see explicitly whether this can leave some room for an SM4+D light darkon if the LHC sees an SM3-like Higgs, we now compare with the corresponding rate in the SM3 without the darkon. Thus for $m_h = 125\,{\rm GeV}$,

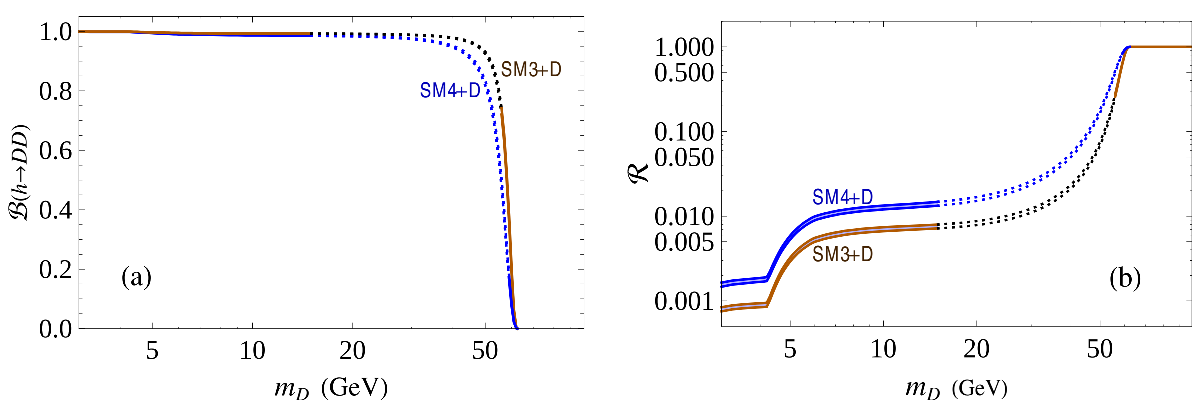

FIG. 2: (a) Branching ratio of $h \to DD$ and (b) the resulting reduction factor $\mathcal{R}$ as functions of m_D in SM3+D and SM4+D for $m_h = 125$ GeV. The dotted parts are disallowed by direct-search limits.

fourth-generation quark masses of order 500 GeV, and $m_D = 15$ GeV, we obtain the factor

$$R_X = \frac{\sigma(pp \to h + \text{anything})_{\rm SM4}\, \mathcal{B}\big(h \to X\bar{X}\big)_{\rm SM4+D}}{\sigma(pp \to h + \text{anything})_{\rm SM3}\, \mathcal{B}\big(h \to X\bar{X}\big)_{\rm SM3}} \sim \frac{9\ \Gamma_h^{\rm SM3}}{\Gamma_h^{\rm SM4} + \Gamma(h \to DD)_{\rm SM4+D}} \tag{3}$$

for $X = \tau^-, c, b, W^{(*)}$, or $Z^{(*)}$ in which case $\Gamma\big(h \to X\bar{X}\big)_{\rm SM4} = \Gamma\big(h \to X\bar{X}\big)_{\rm SM3}$, but for $X = \gamma$ there is extra suppression from $\Gamma(h \to \gamma\gamma)_{\rm SM4} < \Gamma(h \to \gamma\gamma)_{\rm SM3}$ due to the new heavy fermions [31]. In view of Fig. 2(b), for $m_D \le 15$ GeV we find

$$R_X < 0.09\ . \tag{4}$$

We conclude that in the SM4+D with a light darkon the Higgs production event rates would not be SM3-like and the light-darkon region would be ruled out by the detection of such a Higgs, as in the SM3+D case.

In contrast, an SM3+D darkon with $m_D \ge m_h/2$ would be in harmony with the discovery of an SM3-like Higgs, as its decay pattern is not modified by the darkon effect at leading order. However, this mass region up to $m_D \sim 80$ GeV is already forbidden by direct-search limits, as Fig. 1(b) indicates. Darkon masses higher than ∼80 GeV are still viable and will be probed by future direct searches [17]. As for the SM4+D in this m_D region, the detection of such a Higgs would also spell trouble, being at odds with the SM4 prediction of considerably amplified Higgs production cross-sections [30].

III. TWO-HIGGS-DOUBLET MODEL PLUS DARKON

In this darkon model, the Higgs sector is the so-called type III of the two-Higgs-doublet model (THDM). The general form of its Yukawa Lagrangian can be expressed as [32]

$$\begin{aligned}\mathcal{L}_{\rm Y} = & -\bar{Q}_{j,L}\big(\lambda_1^{\mathcal{U}}\big)_{jl}\tilde{H}_1\mathcal{U}_{l,R} - \bar{Q}_{j,L}\big(\lambda_1^{\mathcal{D}}\big)_{jl}H_1\mathcal{D}_{l,R} - \bar{Q}_{j,L}\big(\lambda_2^{\mathcal{U}}\big)_{jl}\tilde{H}_2\mathcal{U}_{l,R} - \bar{Q}_{j,L}\big(\lambda_2^{\mathcal{D}}\big)_{jl}H_2\mathcal{D}_{l,R} \\ & - \bar{L}_{j,L}\big(\lambda_1^{E}\big)_{jl}H_1E_{l,R} - \bar{L}_{j,L}\big(\lambda_2^{E}\big)_{jl}H_2E_{l,R} + \text{H.c.}\ ,\end{aligned} \tag{5}$$

where summation over $j, l = 1, 2, 3$ is implied, $Q_{j,L}$ $\big(L_{l,L}\big)$ denote the left-handed quark (lepton) doublets of the three families, $\mathcal{U}_{l,R}$ and $\mathcal{D}_{l,R}$ $\big(E_{l,R}\big)$ are the right-handed quark (charged lepton) fields, $H_{1,2}$ represent the Higgs doublets, $\tilde{H}_{1,2} = i\tau_2 H_{1,2}^*$, and hence $\lambda_{1,2}^{\mathcal{U},\mathcal{D},E}$ are 3×3 matrices containing the Yukawa couplings. In terms of the Higgs components,

$$H_a = \begin{pmatrix} h_a^+ \\ \frac{1}{\sqrt{2}}\big(v_a + h_a^0 + iI_a^0\big) \end{pmatrix}, \tag{6}$$

where $a = 1, 2$ and v_a is the VEV of H_a satisfying $v_1^2 + v_2^2 = v^2$, with $v = 246$ GeV. The fields h_a^+ and I_a can be expressed in terms of physical Higgs bosons H^+ and A and the would-be Goldstone bosons w and z as

$$\begin{aligned}\begin{pmatrix} h_1^+ \\ h_2^+ \end{pmatrix} &= \begin{pmatrix} \cos\beta & -\sin\beta \\ \sin\beta & \cos\beta \end{pmatrix}\begin{pmatrix} w^+ \\ H^+ \end{pmatrix}, \\ \begin{pmatrix} I_1 \\ I_2 \end{pmatrix} &= \begin{pmatrix} \cos\beta & -\sin\beta \\ \sin\beta & \cos\beta \end{pmatrix}\begin{pmatrix} z \\ A \end{pmatrix},\end{aligned} \tag{7}$$

with $\cos\beta = v_1/v$ and $\sin\beta = v_2/v$, whereas $h^0_{1,2}$ are related to the CP-even Higgs mass eigenstates H and h by

$$\begin{pmatrix} h_1^0 \\ h_2^0 \end{pmatrix} = \begin{pmatrix} \cos\alpha & -\sin\alpha \\ \sin\alpha & \cos\alpha \end{pmatrix} \begin{pmatrix} H \\ h \end{pmatrix} . \tag{8}$$

After the diagonalization of the fermion mass matrices $M_{\mathcal{U},\mathcal{D},E} = \big(\lambda_1^{\mathcal{U},\mathcal{D},E} v_1 + \lambda_2^{\mathcal{U},\mathcal{D},E} v_2\big)/\sqrt{2}$, one can derive from $\mathcal{L}_{\rm Y}$ the Lagrangian for the couplings of $h^0_{1,2}$ to the fermions

$$\begin{aligned} \mathcal{L}'_{\rm Y} \;=\; & -\bar{\mathcal{U}}_L \left[\left(M_{\mathcal{U}} - \frac{\lambda_2^{\mathcal{U}} v_2}{\sqrt{2}} \right) \frac{h_1^0}{v_1} + \left(M_{\mathcal{U}} - \frac{\lambda_1^{\mathcal{U}} v_1}{\sqrt{2}} \right) \frac{h_2^0}{v_2} \right] \mathcal{U}_R \\ & - \bar{\mathcal{D}}_L \left[\left(M_{\mathcal{D}} - \frac{\lambda_2^{\mathcal{D}} v_2}{\sqrt{2}} \right) \frac{h_1^0}{v_1} + \left(M_{\mathcal{D}} - \frac{\lambda_1^{\mathcal{D}} v_1}{\sqrt{2}} \right) \frac{h_2^0}{v_2} \right] \mathcal{D}_R \\ & - \bar{E}_L \left[\left(M_E - \frac{\lambda_2^{E} v_2}{\sqrt{2}} \right) \frac{h_1^0}{v_1} + \left(M_E - \frac{\lambda_1^{E} v_1}{\sqrt{2}} \right) \frac{h_2^0}{v_2} \right] \mathcal{D}_R \;+\; \text{H.c.} \;, \end{aligned} \tag{9}$$

where now $M_{\mathcal{U}} = \text{diag}\big(m_u, m_c, m_t\big)$, etc., and all the fermions in $\mathcal{U} = (u \;\; c \;\; t)^{\rm T}$, etc., are mass eigenstates, but $\lambda_{1,2}^{\mathcal{U},\mathcal{D},E}$ in general are not also diagonal separately. For each of the flavor-diagonal couplings in $\mathcal{L}'_{\rm Y}$, one can then write in terms of the physical field $\mathcal{H} = h$ or H

$$\mathcal{L}_{ff\mathcal{H}} \;=\; -k_f^{\mathcal{H}}\, m_f\, \bar{f} f\, \frac{\mathcal{H}}{v} \;, \tag{10}$$

where for, say, the first family

$$\begin{aligned} k_u^h &= \frac{\cos\alpha}{\sin\beta} - \frac{\lambda_1^u\, v\, \cos(\alpha-\beta)}{\sqrt{2}\, m_u\, \sin\beta} \;, & k_u^H &= \frac{\sin\alpha}{\sin\beta} - \frac{\lambda_1^u\, v\, \sin(\alpha-\beta)}{\sqrt{2}\, m_u\, \sin\beta} \;, \\ k_d^h &= -\frac{\sin\alpha}{\cos\beta} + \frac{\lambda_2^d\, v\, \cos(\alpha-\beta)}{\sqrt{2}\, m_d\, \cos\beta} \;, & k_d^H &= \frac{\cos\alpha}{\cos\beta} + \frac{\lambda_2^d\, v\, \sin(\alpha-\beta)}{\sqrt{2}\, m_d\, \cos\beta} \;, \\ k_e^h &= -\frac{\sin\alpha}{\cos\beta} + \frac{\lambda_2^e\, v\, \cos(\alpha-\beta)}{\sqrt{2}\, m_e\, \cos\beta} \;, & k_e^H &= \frac{\cos\alpha}{\cos\beta} + \frac{\lambda_2^e\, v\, \sin(\alpha-\beta)}{\sqrt{2}\, m_e\, \cos\beta} \;, \end{aligned} \tag{11}$$

where $\lambda_a^{u,d,e} = \big(\lambda_a^{\mathcal{U},\mathcal{D},E}\big)_{11}$. The corresponding $k_f^{\mathcal{H}}$ for the second and third families have analogous expressions. Since only the combination $\lambda_1^f v_1 + \lambda_2^f v_2 = \sqrt{2}\, m_f$ is fixed by the f mass, λ_a^f in $k_f^{\mathcal{H}}$ is a free parameter, and so is $k_f^{\mathcal{H}}$. We remark that setting $\lambda_1^{\mathcal{U}} = \lambda_2^{\mathcal{D}} = \lambda_2^E = 0$ leads to the type II of the THDM+D studied in Refs. [19, 22].

Since the matrices $\lambda_{1,2}^{\mathcal{U},\mathcal{D},E}$ in Eq. (9) generally are not diagonal, their off-diagonal elements may give rise to flavor-changing neutral currents (FCNC) involving the Higgses at tree level. We assume that these flavor-changing elements have their naturally small values according to the Cheng-Sher *ansatz* [33], namely, $(\lambda_a)_{jl} \sim \big(m_j m_l\big)^{1/2}/v$ for $j \neq l$. Since this *ansatz* is facing challenges from current experiments [34], we could suppress the FCNC effects further by increasing the mediating Higgs masses, beyond which fine-tuning may be unavoidable.

Turning to the DM sector of the THDM+D, as in the SM+D, we ensure the darkon's stability as a WIMP candidate by assuming D to be a gauge singlet and introducing a discrete Z_2

symmetry under which $D \to -D$, all the other fields being unaltered. Its renormalizable Lagrangian then takes the form [20, 21]

$$\mathcal{L}_D = \tfrac{1}{2}\partial^\mu D\,\partial_\mu D - \tfrac{1}{4}\lambda_D D^4 - \tfrac{1}{2}m_0^2 D^2 - \left[\lambda_1 H_1^\dagger H_1 + \lambda_2 H_2^\dagger H_2 + \lambda_3\left(H_1^\dagger H_2 + H_2^\dagger H_1\right)\right] D^2 \,. \quad (12)$$

The parameters in the potential of the model should be chosen such that D does not develop a VEV and the Z_2 symmetry stays unbroken, so that D does not mix with the Higgs fields, maintaining its stability. After electroweak symmetry breaking, Eq. (12) contains the darkon mass m_D and the $DD(h,H)$ terms $-\lambda_h v\,D^2 h - \lambda_H v\,D^2 H$, but no DDA coupling, where

$$m_D^2 = m_0^2 + \left[\lambda_1\cos^2\!\beta + \lambda_2\sin^2\!\beta + \lambda_3\sin(2\beta)\right]v^2 \,, \quad (13)$$

$$\begin{aligned} \lambda_h &= -\lambda_1\sin\alpha\,\cos\beta + \lambda_2\cos\alpha\,\sin\beta + \lambda_3\cos(\alpha+\beta) \,, \\ \lambda_H &= \lambda_1\cos\alpha\,\cos\beta + \lambda_2\sin\alpha\,\sin\beta + \lambda_3\sin(\alpha+\beta) \,. \end{aligned} \quad (14)$$

Since m_0 and $\lambda_{1,2,3}$ are free parameters, so are m_D and $\lambda_{h,H}$. We note that for a heavy darkon with $m_D > m_{h,H,A,H^+}$ the darkon annihilation rate also gets contributions from DD couplings to Higgs pairs $\left(h^2, H^2, hH, AA, H^+H^-\right)$ which can be easily derived from Eq. (12).

To evaluate the annihilation rates, the h and H couplings to the W and Z bosons may be relevant depending on m_D. The couplings are given by

$$\mathcal{L}_{VV\mathcal{H}} = \frac{1}{v}\left(2m_W^2\,W^{+\mu}W_\mu^- + m_Z^2\,Z^\mu Z_\mu\right)\left[h\,\sin(\beta-\alpha) + H\,\cos(\beta-\alpha)\right] \quad (15)$$

from the Higgs kinetic sector of the model [32].

Inspired by the possible evidence for a 125-GeV SM-like Higgs in the LHC data, we adopt

$$\cos(\beta-\alpha) = 0 \,. \quad (16)$$

Applying one of its solutions, $\beta-\alpha = \pi/2$, to Eqs. (11), (14), and (15) yields

$$k_u^h = k_d^h = k_e^h = 1 \,, \quad (17)$$

$$k_u^H = -\cot\beta + \frac{\lambda_1^u\,v}{\sqrt{2}\,m_u\,\sin\beta} \,, \qquad k_d^H = \tan\beta - \frac{\lambda_2^d\,v}{\sqrt{2}\,m_d\,\cos\beta} \,,$$
$$k_e^H = \tan\beta - \frac{\lambda_2^e\,v}{\sqrt{2}\,m_e\,\cos\beta} \,, \quad (18)$$

$$\lambda_h = \lambda_1\cos^2\beta + \lambda_2\sin^2\beta + \lambda_3\sin(2\beta) \,, \qquad \lambda_H = \tfrac{1}{2}\left(\lambda_1-\lambda_2\right)\sin(2\beta) - \lambda_3\cos(2\beta) \,, \quad (19)$$

$$\mathcal{L}_{VV\mathcal{H}} = \left(2m_W^2\,W^{+\mu}W_\mu^- + m_Z^2\,Z^\mu Z_\mu\right)\frac{h}{v} \,. \quad (20)$$

Another consequence is that the tree-level flavor-changing couplings of h vanish. Evidently, now the couplings of h to the SM fermions and gauge bosons are identical to those of the SM Higgs. The alternative solution, $\beta-\alpha = -\pi/2$, would yield the same results, but with the opposite signs. It is worth remarking that Eq. (16) is part of the $0 \le |\cos(\beta-\alpha)| \ll 1$ region of the model parameter space where h can have SM-like couplings to the fermions and gauge bosons,

provided that $\tan\beta$ and $\cot\beta$, as well as the Higgs self-couplings, are not large and that the A mass is not below the electroweak scale [35].

To render h more SM-like, we require that the hDD coupling $\lambda_h = 0$.[4] It follows that for $m_D < m_h < m_H$ the darkon annihilation contribution to the DM relic density comes only from H-mediated diagrams. Since λ_1^u and $\lambda_2^{d,e}$ in Eq. (18) are free parameters, for illustration we will pick for definiteness

$$k_u^H = k_d^H = k_e^H = 1 , \tag{21}$$

and similarly for k_f^H belonging to the second and third families. With these specific selections, H share with h the same couplings to the fermions, but H does not couple to the W and Z bosons at tree level, unlike h.

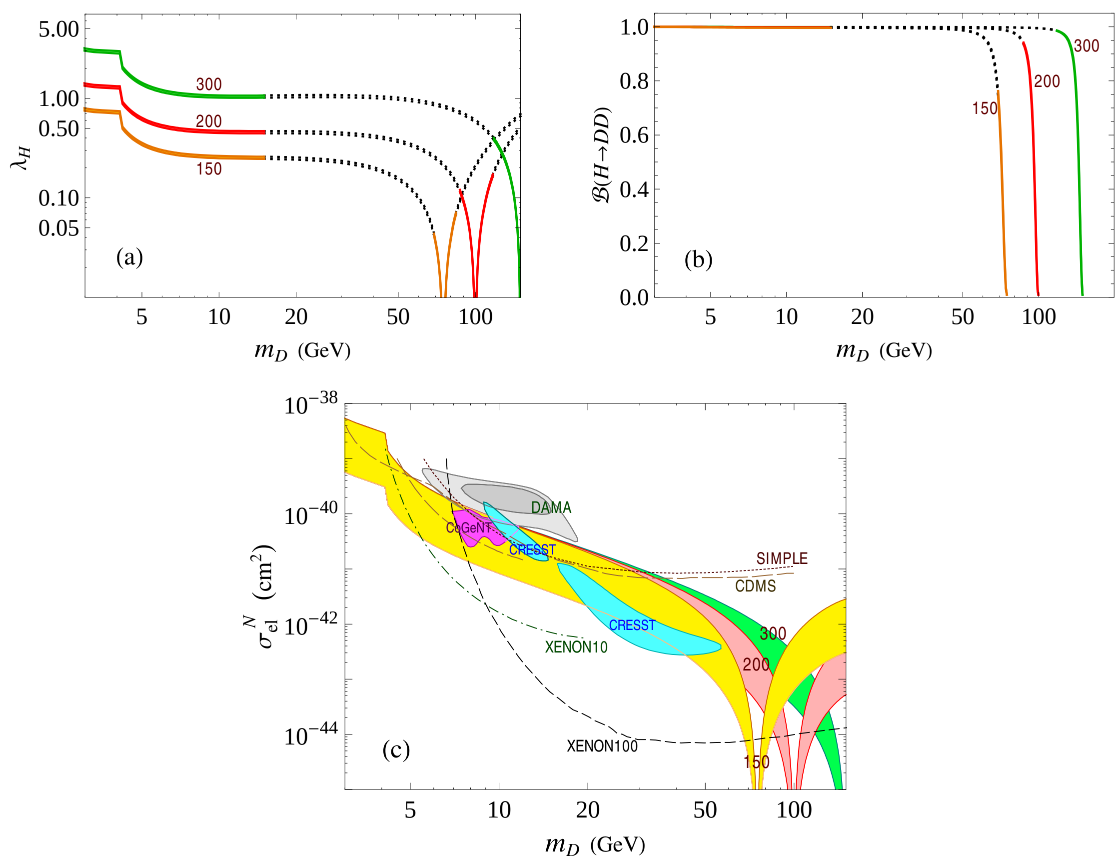

FIG. 3: (a) Darkon-H coupling λ_H as a function of darkon mass m_D for $m_H = 150, 200, 300$ GeV, with the other couplings specified in the text, in the THDM+D with isospin-conserving darkon-nucleon interactions. The resulting (b) branching ratio of $H \to DD$ and (c) darkon-nucleon cross-section $\sigma_{\rm el}^N$, compared to the same experimental results as in Fig. 1(b). The black-dotted parts in (a) and (b) are disallowed by direct-search bounds.

[4] The value of λ_h may deviate somewhat from zero if h is found with an invisible decay rate that exceeds its SM range. Since $\lambda_{1,2,3}$ in Eq. (19) are free parameters, the determination of λ_h does not fix λ_H.

Since we are here interested in the case of a relatively light darkon, we concentrate on the $m_D \leq 150\,\text{GeV}$ region. Upon specifying m_H, one can extract λ_H from the relic-density data, the procedure being similar to that in the preceding SM+D case. We present the results for some illustrative values of m_H in Fig. 3(a), where as before the dotted regions are forbidden by direct-search data. From now on, we assume $m_D \leq m_H < m_{A,H^\pm}$.

The extracted λ_H translates into the branching ratio of invisible decay $H \to DD$ in Fig. 3(b). As expected in the absence of $H(WW,ZZ)$ couplings at tree level, $\mathcal{B}(H \to DD)$ stays close to 1 over most of the kinematically permitted range. It is different from $\mathcal{B}(h \to DD)$ in the SM+D which becomes significantly less dominant if $m_h > 2m_{W,Z}$ after the important $h \to WW, ZZ$ channels are open [17, 18].

The cross section of elastic scattering of a darkon off a nucleon N mediated by H is [19]

$$\sigma_{\rm el}^N \;=\; \frac{\lambda_H^2\, g_{NNH}^2\, v^2\, m_N^2}{\pi \left(m_D + m_N\right)^2 m_H^4} \;, \tag{22}$$

where g_{NNH} is the effective H-nucleon coupling. With the choices $k_f^H = 1$, as in Eq. (21), which conserve isospin, g_{NNH} is no different from g_{NNh} used in the previous section. Hence we employ $0.0011 \leq g_{NNH} \leq 0.0032$.

We show in Fig. 3(c) the calculated $\sigma_{\rm el}^N$ corresponding to the parameter selections in Fig. 3(a). Also on display are the results of recent DM direct searches, as in Fig. 1(b). One can see that, much like the SM+D case, the THDM+D prediction for the darkon-nucleon cross-section can accommodate well especially the light-WIMP regions suggested by CoGeNT and CRESST-II data [4, 5], although they are in tension with the null results of other direct searches. However, unlike the SM+D, the THDM+D still has enough parameter space to allow the particular choices we made which yield both a SM-like Higgs, h, and a DM sector possessing a viable light WIMP coupled to another Higgs, H.

IV. ISOSPIN-VIOLATING DARK MATTER IN THDM+D

Up to now, we have only achieved making the THDM+D have a low-mass WIMP candidate and a SM-like Higgs boson as hinted at by the LHC findings. If the model is also to accommodate both DM direct-search results which indicated light-WIMP evidence and those which did not, it needs to have a mechanism that can provide substantial isospin violation in the WIMP interactions with nucleons. As was proposed in the literature [12–14], the tension in the light-WIMP data will partially go away if the WIMP effective couplings $f_{p,n}$ to the proton and neutron, respectively, satisfy the relation $f_n \sim -0.7 f_p$. The THDM+D may be able to realize this using the freedom still available in the parameters k_f^H defined above. By allowing them to deviate from the choices $k_f^H = 1$ in the last section, which respect isospin, it may be feasible for the model to attain the desired results. We explore this scenario in the following.

The WIMP-nucleon cross-section $\sigma_{\rm el}^N$ in the isospin-symmetric limit can be expressed in terms of the WIMP-proton elastic cross-section $\sigma_{\rm el}^p$ in the presence of isospin violation as [12, 13]

$$\sigma_{\rm el}^N\, f_p^2 \sum_i \eta_i\, \mu_{A_i}^2\, A_i^2 \;=\; \sigma_{\rm el}^p \sum_i \eta_i\, \mu_{A_i}^2 \big[\mathcal{Z} f_p + \big(A_i - \mathcal{Z}\big) f_n\big]^2 \;, \tag{23}$$

where the sum is over the isotopes of the element in the detector material with which the WIMP interacts dominantly, $\mathcal{Z}$ is proton number of the element, A_i (η_i) each denote the nucleon number

(fractional abundance) of its isotopes, and $\mu_{A_i} = m_{A_i} m_{\rm WIMP}/\big(m_{A_i}+m_{\rm WIMP}\big)$ involving the isotope and WIMP masses. Thus if isospin violation is negligible, $f_n = f_p$, the measurement of event rates of WIMP-nucleus scattering will translate into the usual $\sigma_{\rm el}^N = \sigma_{\rm el}^p$. For $f_n = -0.7 f_p$, taking into account the different A_i and $\mathcal{Z}$ numbers for the different detector materials, one can transform some of the contradictory data on the WIMP-nucleon cross-sections into $\sigma_{\rm el}^p$ numbers which overlap with each other [13, 14]. This also makes the extracted $\sigma_{\rm el}^p$ enhanced relative to the current measured values of $\sigma_{\rm el}^N$ by up to 4 orders of magnitude, depending on A_i and $\mathcal{Z}$.

Now, in the THDM+D with only H mediating the WIMP-nucleon interactions

$$\sigma_{\rm el}^p \;=\; \frac{4\, m_D^2\, m_p^2\, f_p^2}{\pi\big(m_D+m_p\big)^2} \;, \qquad f_p \;=\; \frac{\lambda_H\, g_{ppH}\, v}{2\, m_D\, m_H^2} \;, \tag{24}$$

where the H-proton effective coupling g_{ppH} contains various quark contributions according to Eq. (A1). The relation $f_n = \rho\, f_p$ then implies

$$g_{nnH} \;=\; \rho\, g_{ppH} \;, \tag{25}$$

where the H-neutron effective coupling g_{nnH} also has the general form in Eq. (A1) and we will set $\rho = -0.7$. Also relevant is the darkon annihilation-rate formula [16, 19]

$$\sigma_{\rm ann} v_{\rm rel} \;=\; \frac{4\lambda_H^2 v^2}{\big(4m_D^2-m_H^2\big)^2+\Gamma_H^2\, m_H^2}\, \frac{\sum_i \Gamma\big(\tilde H\to X_i\big)}{m_D} \;, \tag{26}$$

where $v_{\rm rel}$ is the relative speed of the DD pair in their center-of-mass frame, $\tilde H$ is a virtual H having the same couplings to other states as the physical H of mass m_H, but with an invariant mass $\sqrt{s} = 2m_D$, and $\tilde H \to X_i$ is any kinematically permitted decay mode of $\tilde H$.

Using Eqs. (24)-(26) and focusing on the low-mass range $5\,{\rm GeV} \le m_D \le 20\,{\rm GeV}$, we scan the parameter space of the products $\lambda_H k_f^H$, as the factors always go together in Eqs. (24) and (26), while imposing $\rho = -0.7$ and the relic-density constraint. We find that to enhance $\sigma_{\rm el}^p$ by a few orders of magnitude under these restrictions implies that $k_{u,d}^H$ have to be big, $k_u^H \sim -2k_d^H$, and the other k_f^H become negligible by comparison, confirming the finding of Ref. [36]. For instance, with $m_H = 200\ (300)\,{\rm GeV}$ we obtain $0.6\ (1.4)\times 10^3 \le \lambda_H k_u^H \le 0.8\ (1.8)\times 10^3$ corresponding to $5\,{\rm GeV} \le m_D \le 20\,{\rm GeV}$. It follows that in general $k_u^H = \mathcal{O}\big(10^3\big)$ if $\lambda_H = \mathcal{O}(1)$ and m_H is a few hundred GeV. For such large $k_{u,d}^H$, one expects that $k_u^H \sim \lambda_1^u v_1/m_u$ and $k_d^H \sim \lambda_2^d v_2/m_d$ from Eq. (18). Consequently, since $\lambda_1^u v_1 + \lambda_2^u v_2 = \sqrt{2}\, m_u$ and $\lambda_1^d v_1 + \lambda_2^d v_2 = \sqrt{2}\, m_d$, some degree of subtle cancelations between the $\lambda_a^{u,d} v_a$ terms is needed to reproduce the small u and d masses. This is the price one has to pay for the greatly amplified $\sigma_{\rm el}^p$.

We plot the theoretical cross-section (orange curve) in Fig. 4, where the band width reflects the relic-density uncertainty. Also plotted are the direct-search results [3–9] reproduced with the WIMP-nucleon couplings satisfying $f_n = -0.7 f_p$. In this m_D range, the theory curve is roughly independent of m_H if it is in the hundreds of GeV. One can see that the prediction is not able to reach the (gray) region implied by DAMA at the 3σ level and only marginally covers its 5σ (lighter gray) region. The CoGeNT preferred (magenta) area is also unreachable. Nevertheless, with respect to many points in these regions, the prediction is too low by no more than a factor of 2 or 3. On the other hand, with appropriately lowered $k_{u,d}^H$, it can agree well with the cross sections (cyan patch) favored by the CRESST-II results. Furthermore, for

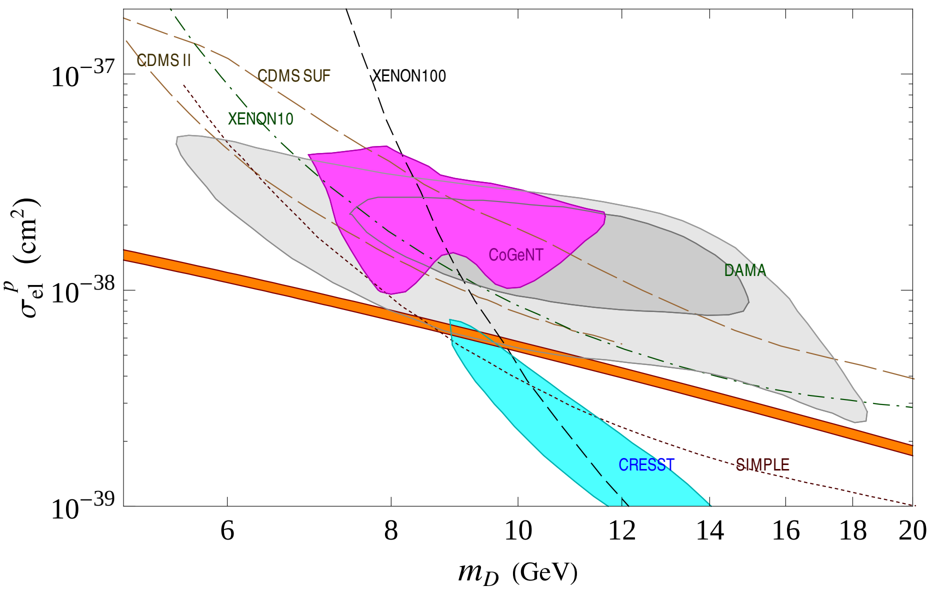

FIG. 4: Darkon-proton cross-section $\sigma_{\rm el}^{p}$ in THDM+D with isospin-violating light darkon (orange curve) compared with several direct-search results for WIMP-nucleon couplings satisfying $f_n = -0.7 f_p$.

m_D values below 10 GeV or so the prediction curve does not conflict with the XENON limits, which is unlike the isospin-symmetric case illustrated in Fig. 3. The situation is very different when it comes to the limit from SIMPLE which now rules out $m_D \gtrsim 9\,\mathrm{GeV}$, while previously it allowed most of the darkon parameter space. On the experimental side, due to $f_n = -0.7 f_p$, virtually none of the CRESST-II area in Fig. 3 is consistent with the DAMA and CoGeNT ones any more, and the SIMPLE bound disallows most of the CoGeNT area. It is obvious from a comparison of Figs. 3(c) and 4 that there are unresolved puzzles remaining. To address them in a comprehensive manner would likely have to await future direct-searches with improved precision and may need to involve extra ingredients [14] beyond the simple frameworks treated in this work.

We now demonstrate that the result above for the enhanced $\sigma_{\rm el}^{p}$ prediction does not depend on the value of k_u^H or the pion-nucleon sigma term $\sigma_{\pi N}$, assuming that $k_{u,d}^H$ are much bigger than the other k_f^H. As discussed in Appendix A, the requirements $k_{u,d}^H \gg k_{s,c,b,t}^H \sim 0$ and $g_{nnH} = \rho\, g_{ppH}$ result in

$$g_{ppH} \;=\; \frac{3.0\times10^{-5}\; k_u^H\,\sigma_{\pi N}}{5.2(1+\rho)\,\mathrm{MeV} \;+\; (1-\rho)\sigma_{\pi N}} \tag{27}$$

and also $k_d^H \propto k_u^H$ according to Eq. (A9). Hence for $-1 \lesssim \rho \lesssim -0.5$ and [18] $\sigma_{\pi N} \geq 30\,\mathrm{MeV}$

$$g_{ppH} \;\sim\; \frac{3\times10^{-5}\; k_u^H}{(1-\rho)}\;, \qquad k_d^H \;\sim\; -0.5\,k_u^H \tag{28}$$

approximately independent of $\sigma_{\pi N}$. Since now g_{ppH}^2 and $\Sigma_i \Gamma_{\tilde H\to X_i} \simeq \Gamma_{\tilde H\to u\bar u} + \Gamma_{\tilde H\to d\bar d}$ are both

roughly proportional to $\left(k_u^H\right)^2$, the approximate expression

$$\sigma_{\rm el}^p \simeq \frac{g_{ppH}^2\, m_D\, m_p^2}{4\pi \left(m_D+m_p\right)^2} \, \frac{\sigma_{\rm ann} v_{\rm rel}}{\Gamma_{\tilde H\to u\bar u}+\Gamma_{\tilde H\to d\bar d}} \,, \tag{29}$$

derived from Eqs. (24) and (26) and valid for $m_p < m_D \ll m_H$, is roughly independent of k_u^H. Consequently, $\sigma_{\rm el}^p$ can only be increased further if $\sigma_{\rm ann} v_{\rm rel}$ is also increased.

Finally, it is worth mentioning that one cannot obtain the substantial isospin violation of interest if the Higgs sector of the THDM+D is of type II, in which the up- and down-type quarks are coupled to different Higgs doublets [32]. In that case, the second term in each of the formulas for k_f^H in Eq. (18) is absent, and therefore there is not much freedom to vary k_f^H, as $\tan\beta$ cannot be arbitrarily small or big due to restrictions from various data [2, 37]. Accordingly, since the darkon-nucleon couplings are dominated by the combined strange- and heavy-quark contributions, which conserve isospin, as can be seen from Eqs. (A1) and (A6), the darkon-proton coupling cannot be made large enough after applying the relic density and $g_{nnH} = -0.7\, g_{ppH}$ restraints. We have found that the resulting darkon-proton cross-section $\sigma_{\rm el}^p$ for $m_D \sim 10\,{\rm GeV}$ cannot reach more than $\sim 10^{-43}\,{\rm cm}^2$, about 5 orders of magnitude too small.

V. CONCLUSIONS

The preliminary indications of a Higgs boson with SM-like properties in the latest LHC data, if confirmed by future measurements, will have important implications for WIMP dark matter models. We have explored a number of such implications for some of the simplest darkon models. For the simplest one, SM3+D, most of the light-darkon mass range will be ruled out if an SM3-like Higgs with mass near 125 GeV is found. Such a discovery would also be at odds with the SM4 Higgs expectations, thus disfavoring the SM4+D. In contrast, the type-III two-Higgs-doublet model enlarged with the addition of a darkon has an abundance of allowed parameter space in its DM sector. It can accommodate an SM3-like Higgs and simultaneously offers a WIMP candidate in harmony with the light-WIMP hypothesis inspired by the clues from a number of DM searches, although it is in conflict with the null results of other searches. The model can also provide substantial isospin violation in the WIMP-nucleon interactions which can alleviate some of this tension in the data on light DM. However, this could be achieved only with some amount of fine-tuning in several of the relevant parameters. Nevertheless, upcoming searches for the Higgs at the LHC and future DM direct searches can test further the THDM+D which we have considered. Additional signals of the darkon in this scenario may be available from flavor-changing top-quark decays into lighter up-type quarks with missing energy which are potentially observable at the LHC.

Acknowledgments

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Appendix A: Higgs-nucleon effective couplings

The effective coupling of a Higgs $\mathcal{H}$ to a proton p or neutron n is related to the quark Yukawa parameters by [19, 38]

$$g_{\mathcal{N}\mathcal{N}\mathcal{H}}\,\bar{u}_{\mathcal{N}}u_{\mathcal{N}} \;=\; \bar{u}_{\mathcal{N}}u_{\mathcal{N}}\sum_q g_q^{\mathcal{N}}k_q^{\mathcal{H}} \;=\; \sum_q \frac{k_q^{\mathcal{H}}}{v}\langle\mathcal{N}|m_q\,\bar{q}q|\mathcal{N}\rangle\;, \qquad \mathcal{N} \;=\; p,n\;, \tag{A1}$$

where $u_{\mathcal{N}}$ is the Dirac spinor for $\mathcal{N}$ and the sum is over all quarks. Using the chiral Lagrangian approach described in Ref. [19], but without neglecting isospin violation, and assuming three fermion families, we obtain

$$g_u^p \;=\; \frac{-2\big(b_D+b_F+b_0\big)m_u}{v}\;, \qquad g_u^n \;=\; \frac{-2b_0\,m_u}{v}\;, \tag{A2}$$

$$g_d^p \;=\; \frac{-2b_0\,m_d}{v}\;, \qquad g_d^n \;=\; \frac{-2\big(b_D+b_F+b_0\big)m_d}{v}\;, \tag{A3}$$

$$g_s^p \;=\; g_s^n \;=\; \frac{-2\big(b_D-b_F+b_0\big)m_s}{v}\;, \qquad g_{c,b,t}^p \;=\; g_{c,b,t}^n \;=\; \frac{2\,m_B}{27\,v}\;, \tag{A4}$$

$$\sigma_{\pi N} \;=\; -\big(b_D+b_F+2b_0\big)\big(m_u+m_d\big)\;, \tag{A5}$$

where the parameters $b_{D,F,0}$ and m_B can be fixed from the measured masses of the lightest baryons and the phenomenological or lattice value of the pion-nucleon sigma term $\sigma_{\pi N}$. Since $\sigma_{\pi N}$ is still poorly determined [28], we take $30\,\mathrm{MeV} \le \sigma_{\pi N} \le 80\,\mathrm{MeV}$ after Ref. [18]. Hence for $\sigma_{\pi N} = 30$ (80) MeV the values of $g_q^{\mathcal{N}}$ are, in units of 10^{-3},

$$\begin{aligned} g_u^p &= 0.05\ (0.12)\;, & g_u^n &= 0.04\ (0.11)\;,\\ g_d^p &= 0.06\ (0.20)\;, & g_d^n &= 0.09\ (0.22)\;,\\ g_s^{p,n} &= 0.25\ (2.88)\;, & g_{c,b,t}^{p,n} &= 0.26\ (0.05). \end{aligned} \tag{A6}$$

In the following we assume that $k_{u,d}^H \neq 0$, the other k_f^H are zero, and $g_{nnH} = \rho\, g_{ppH}$ as in Eq. (25), with ρ being a constant. Combining these requirements with some of the preceding equations, we arrive at

$$k_d^H \;=\; \frac{(1+\rho)\big(b_D+b_F\big)\big(m_u+m_d\big)+(1-\rho)\sigma_{\pi N}}{(1+\rho)\big(b_D+b_F\big)\big(m_u+m_d\big)-(1-\rho)\sigma_{\pi N}}\;\frac{m_u}{m_d}\,k_u^H\;, \tag{A7}$$

$$g_{ppH} \;=\; \frac{4\big(b_D+b_F\big)m_u\,k_u^H\,\sigma_{\pi N}/v}{\big(b_D+b_F\big)\big(m_u+m_d\big)(1+\rho)-(1-\rho)\sigma_{\pi N}}\;. \tag{A8}$$

Numerically we get $\big(b_D+b_F\big)\big(m_u,m_d\big) \simeq (-1.9,-3.3)\,\mathrm{MeV}$, which in this case leads to

$$k_d^H \;\simeq\; \frac{2.9(1+\rho)\,\mathrm{MeV}\;-\;0.56\,(1-\rho)\sigma_{\pi N}}{5.2(1+\rho)\,\mathrm{MeV}\;+(1-\rho)\sigma_{\pi N}}\;k_u^H\;, \tag{A9}$$

$$g_{ppH} \;\simeq\; \frac{3.0\times10^{-5}\;k_u^H\,\sigma_{\pi N}}{5.2(1+\rho)\,\mathrm{MeV}\;+\;(1-\rho)\sigma_{\pi N}}\;. \tag{A10}$$

In evaluating these quantities and $\Sigma_i \Gamma_{\tilde{H}\to X_i}$ in Eq. (26), we have employed the running masses of the light quarks at scales $\mu = 1\,\mathrm{GeV}$ and $\mu = 2m_D$, respectively, and included QCD corrections in the $\tilde{H} \to q\bar{q}$ rates [39]. The ratios of light-quark masses used are $m_d/m_u \simeq 1.8$ and $m_s/m_d \simeq 20$, which fall within their measured ranges [40].

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